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Chaotic inflation and a radiatively generated intermediate scale in the supersymmetric standard model

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Abstract

We consider a phenomenological extension of the minimal supersymmetric standard model which incorporates chaotic inflation and a radiatively generated intermediate mass scale. Initially a period of chaotic inflation is driven by a quartic potential associated with the right-handed electron sneutrino. Supersymmetry relates the quartic coupling of the inflationary potential to the electron Majorana neutrino Yukawa coupling, h_1. The microwave background temperature anisotropy determines this coupling to be $h_1 \simeq 10^{-7}$, which is similar in magnitude to the electron Dirac Yukawa coupling. A U(1) Peccei-Quinn (PQ) symmetry is broken by radiative corrections at an intermediate scale $\simeq 10^{12}$GeV when the universe cools to a temperature $T \lesssim 10^3$GeV. This leads to an invisible axion, a weak scale μ-term and an electron Majorana neutrino mass $M_{N_1} \simeq 10^5$GeV. A second inflationary period can also occur via a flat-direction field. In this case the universe can be reheated to a temperature $T_{RH} \simeq 10^6$GeV, without restoring PQ symmetry. Baryogenesis will then occur via out-of-equilibrium neutrino decay.

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1 Introduction

While the minimal supersymmetric standard model (MSSM) provides the most promising extension of the successful standard model, it does not yet encompass important ideas that would be expected of the complete low energy effective theory. These include neutrino masses, baryogenesis, inflation and the strong CP problem. Recently there have been a number of interesting proposals which partly address these shortcomings. In particular several papers by Murayama, Yanagida and collaborators have made significant contributions [1]; see also ref. [2]. However a more complete phenomenological model is lacking at present.

In this work we construct a phenomenological extension of the MSSM which can successfully incorporate inflation, neutrino masses, baryogenesis and axions. We build on the work of Murayama et al [1, 3], modifying their approach in a way specified below. The resulting Lagrangian can describe all of the usual supersymmetry phenomenology, including cold dark matter, LEP data and BR($b \to s\gamma$). In addition it includes neutrino masses via a see-saw mechanism (which provide the hot dark matter), induces inflation with a sneutrino inflaton, incorporates baryogenesis and accommodates the axion. Interestingly the parameters relevant to each of these ideas are closely interrelated in our model. Furthermore, the fine-tuning in the inflationary potential is no more worse than that of the electron Yukawa coupling in the MSSM. We will postpone giving any detailed calculations here, and instead outline how such an encompassing Lagrangian can be constructed.

In order to incorporate chaotic inflation one needs a scalar field in the theory to have an initial value much greater than M_{Planck} in the early universe. An unnatural fine-tuning required for gauge non-singlet fields along D-term flat directions restricts the inflaton to be a gauge singlet field such as a sneutrino. This idea of identifying the right-handed sneutrino as the inflaton is due to Murayama et al [1], where it was noted that the addition of the superpotential term $W = \frac{1}{2} M \hat{N}_i^c \hat{N}_i^c$ with a common Majorana mass $M \simeq 10^{13}$GeV coincides with a successful implementation of chaotic inflation using a quadratic scalar potential. However to solve the strong CP problem one expects the Lagrangian in the early universe to be Peccei-Quinn (PQ) invariant. A PQ-invariant Majorana term for the right-handed neutrinos can be written by introducing a singlet superfield $\hat{P}$ with superpotential $W = \frac{1}{2} h_i \hat{N}_i^c \hat{N}_i^c \hat{P}$. This means that in the early universe the inflationary potential will be quartic with a coupling h_1^2. Anisotropic temperature fluctuations, $\delta T/T \simeq 10^{-5}$ in the present universe then determine the Majorana Yukawa

coupling $h_1 \simeq 10^{-7}$ [4].

Neutrino masses via a see-saw mechanism will be generated when the scalar component $\tilde{P}$ of the superfield $\hat{P}$ receives a vacuum expectation value at an intermediate scale. Intermediate scale breaking in the supersymmetric standard model was previously considered by Murayama, Suzuki and Yanagida [3]. Radiative corrections from right-handed neutrino loops break $\mathrm{U}(1)_{PQ}$ by driving the squared mass of $\tilde{P}$ negative at an intermediate scale. This is similar to the normal radiative electroweak symmetry breaking induced in the MSSM by the large top Yukawa coupling. It turns out that a second singlet superfield $\hat{P}'$ is also needed to ensure an invisible axion. The PQ symmetry can only be broken after inflation ends because during the inflationary epoch the inflaton induces an effective $\tilde{P}$ mass, which dominates the radiative corrections from neutrino loops. As the inflaton undergoes coherent oscillations about its minimum, the oscillation amplitude falls off as R^{-1} (R is the scale factor of the universe) for a quartic potential. However as the universe is reheated, finite temperature corrections induce a local minimum at $\langle\tilde{P}\rangle = 0$ and the field $\tilde{P}$ can remain trapped there until $T \lesssim 10^3$GeV.

If a second period of inflation were to commence when the sneutrino is oscillating to zero, the universe would then be supercooled below $T \simeq 10^3$GeV. The potential barrier at $\langle\tilde{P}\rangle = 0$ would disappear and $\tilde{P}$ would drop to the true vacuum at $\langle\tilde{P}\rangle \simeq 10^{12}$GeV. This second inflationary epoch can be caused by a scalar field with amplitude $\mathcal{O}(M_{Pl})$. Typically these scalar fields only have non-renormalisable inflaton couplings and are associated with a flat direction of the supersymmetric theory. Their amplitudes can be driven to values $\mathcal{O}(M_{Pl})$ during the first inflationary epoch [5]. Thus it is likely that after the first inflation period is over there exists some flat-direction field, η with an amplitude $\mathcal{O}(M_{Pl})$ which starts the second inflationary epoch. In contrast to the initial period of chaotic inflation where $V(\phi) \lesssim M_{Pl}^4$, the potential along the flat direction is $V \simeq m_W^2 M_{Pl}^2 \simeq (10^{11}\mathrm{GeV})^4$, where $m_W \simeq \mathcal{O}(\mathrm{TeV})$. This has been referred to as 'intermediate scale inflation' [6] and conveniently coincides with the $\mathrm{U}(1)_{PQ}$ symmetry breaking.

When the second inflationary epoch ends, the universe is reheated to a temperature $T_{RH} \simeq 10^6$GeV which is low enough to prevent restoring PQ symmetry (at the local minimum $\langle\tilde{P}\rangle = 0$). This reheat temperature is high enough for baryogenesis to occur via the out-of-equilibrium decay of the light electron Majorana neutrino (N_1). The initial chaotic inflationary epoch with the right-handed electron sneutrino inflaton solves the flatness and horizon problems and generates the required density perturbations. In order not to wipe out the density perturbations from the original inflationary epoch, we require

that the number of e-foldings, N during the second period of inflation satisfy $N \lesssim 30$ [7]. Note, however that the axion strings resulting from the spontaneous symmetry breakdown will not be completely diluted during the second inflationary epoch.

The main points of our model which differ from previous attempts are as follows. Initially chaotic inflation occurs with a quartic potential associated with the right-handed electron sneutrino. COBE data on the temperature anisotropy then determine the electron Majorana Yukawa coupling to be $h_1 \simeq 10^{-7}$ which is no less fine-tuned than the electron Dirac Yukawa coupling. When the universe is reheated, finite temperature corrections induce a local minimum at $\langle \tilde{P} \rangle = 0$ which persists until $T \simeq 10^3$GeV. If instead a second inflationary epoch occurs at an intermediate scale (via a flat-direction field η), the universe will be supercooled below $T \simeq 10^3$GeV as $\langle \tilde{N}_1^c \rangle \to 0$. Soft breaking terms then dominate and radiatively generate an intermediate mass scale $\simeq 10^{12}$GeV. The mass of the electron Majorana neutrino will typically be $M_{N_1} \simeq 10^5$GeV (rather than the more common value $M_{N_1} \simeq 10^{11}$GeV). When the flat-direction field decays it can reheat the universe to a temperature $T_{RH} \simeq 10^6$ GeV. All supersymmetry breaking effects are parameterised by soft terms in the scalar potential and we do not consider any effects that might arise from supergravity or string theory. The details of our scenario are presented below.

2 Chaotic inflation in the supersymmetric standard model

Consider a PQ invariant extension of the MSSM which provides the framework for our model of inflation. This extension was first considered by Murayama, Suzuki and Yanagida [3]. If a right-handed neutrino field $\hat{N}^c$ is introduced into the MSSM, the possible Yukawa couplings in the superpotential are

$$W[\Phi] = h_U^{ij} \hat{u}_i^c \hat{Q}_j \hat{H}_u + h_D^{ij} \hat{d}_i^c \hat{Q}_j \hat{H}_d + h_E^{ij} \hat{e}_i^c \hat{L}_j \hat{H}_d + h_N^{ij} \hat{N}_i^c \hat{L}_j \hat{H}_u \tag{1}$$

where $\hat{Q}$, $\hat{L}$ and $\hat{H}_{u,d}$ are SU(2) doublet chiral superfields and $\hat{u}^c$,$\hat{d}^c$,$\hat{e}^c$ and $\hat{N}^c$ are SU(2) singlet chiral superfields. The labels i, j are generation indices and all group indices have been suppressed. Notice that to generate a Majorana mass term for the right-handed neutrino only requires coupling $\hat{N}^c$ to a singlet superfield $\hat{P}$. However, with just the superfield $\hat{P}$, the PQ symmetry is broken at the electroweak scale and gives rise to a standard visible axion which has been ruled out experimentally. This problem is avoided by introducing a second singlet field, $\hat{P}'$ which causes the PQ symmetry to be broken

at an intermediate scale and leads to an invisible axion [3]. Thus the most general PQ invariant superpotential with an intermediate breaking scale is given by

$$W'[\Phi] = \tfrac{1}{2} h_M^{ij} \hat{N}_i^c \hat{N}_j^c \hat{P} + \frac{f}{M_{Pl}} \hat{P}^3 \hat{P}' + \frac{g}{M_{Pl}} \hat{P} \hat{P}' \hat{H}_u \hat{H}_d, \tag{2}$$

where M_{Pl} is the Planck mass and the PQ charge assignments are $+1/2$ for $\hat{Q}$, $\hat{L}$, $\hat{u}^c$, $\hat{d}^c$, $\hat{e}^c$, $\hat{N}^c$, -1 for $\hat{P}$, $\hat{H}_{u,d}$ and $+3$ for $\hat{P}'$. The total superpotential of our phenomenological model is $W + W'$. Notice that by introducing $\hat{P}'$ one naturally generates a coupling to the Higgs superfields, which ultimately becomes the μ-term of the MSSM.

For chaotic inflation to occur one needs an inflationary potential $V(\phi) \lesssim M_{Pl}^4$ and a scalar field with an initial value $\phi(0) \gg M_{Pl}$ [8]. The scalar potential resulting from $W + W'$ restricts the amplitude of any gauge non-singlet scalar fields to be $\mathcal{O}(M_{Pl})$ because of unnatural fine-tunings along D-term flat directions [1].[3] This leaves only the scalar components $\tilde{N}_i^c$, $\tilde{P}$ and $\tilde{P}'$ of the singlet superfields $\hat{N}_i^c$,$\hat{P}$ and $\hat{P}'$ as possible candidates for the inflaton. The scalar potential arising from the superpotential W' is given by

$$\begin{aligned} V(\phi) &= \left| \tfrac{1}{2} h_i \tilde{N}_i^c \tilde{N}_i^c + 3 \frac{f}{M_{Pl}} \tilde{P}^2 \tilde{P}' + \frac{g}{M_{Pl}} H_u H_d \tilde{P}' \right|^2 + \left| \tilde{P} \right|^2 \left| \frac{f}{M_{Pl}} \tilde{P}^2 + \frac{g}{M_{Pl}} H_u H_d \right|^2 \\ &+ h_i^2 \left| \tilde{N}_i^c \right|^2 \left| \tilde{P} \right|^2 + \frac{g^2}{M_{Pl}^2} \left| \tilde{P} \right|^2 \left| \tilde{P}' \right|^2 (H_u^\dagger H_u + H_d^\dagger H_d) \end{aligned} \tag{3}$$

where we have assumed for simplicity that the Majorana Yukawa couplings are real and diagonal, $h_M^{ij} = h_i \delta^{ij}$ (the soft breaking terms are not important in this initial inflationary stage and will be considered later). If the couplings $f, g \sim 0.01$ [4] then the condition $V(\phi) \lesssim M_{Pl}^4$ restricts the amplitudes of $\tilde{P}$ and $\tilde{P}'$ to be $\mathcal{O}(M_{Pl})$ which is not enough to solve the flatness and horizon problems. This leaves the right-handed sneutrino as the only candidate for the inflaton. Clearly the lightest sneutrino will end up being the inflaton because it is assumed to have the flatest potential (or smallest Majorana Yukawa coupling). The heavier generations will roll to their minimum fairly quickly because their potential is steeper. In addition the $\tilde{P}$ and $\tilde{P}'$ scalar fields receive induced masses of $\mathcal{O}(M_{Pl})$ and are also driven to their minima early on. Thus if we suppose that the right-handed electron sneutrino acts as the inflaton with $\tilde{N}_1^c(0) \gg M_{Pl}$ then during inflation the potential (3) effectively becomes

$$V(\phi) = \frac{1}{4} h_1^2 \left| \tilde{N}_1^c \right|^4, \tag{4}$$

[3] Also a flat inflationary potential is much more difficult to achieve with gauge couplings.
[4] We will show later that intermediate scale breaking requires $f \gtrsim 0.01$.

with $\langle\tilde{N}^c_{2,3}\rangle, \langle\tilde{P}\rangle, \langle\tilde{P}'\rangle \ll M_{Pl}$. Note that the Higgs ($H_u$) and slepton scalar fields receive induced masses from the inflaton, $\langle\tilde{N}^c_1(t)\rangle$, which are bigger than the Hubble constant H. Consequently the $\mathcal{O}(M_{Pl})$ amplitudes of these fields will be damped away exponentially during the inflationary period.

The inflationary potential (4) is known to generate the required density perturbations for large scale structure of the universe, provided that the quartic coupling (h_1^2) is approximately 10^{-14} [9]. This means that the Majorana Yukawa coupling for the right-handed electron neutrino must be $h_1 \simeq 10^{-7}$. This is similar in magnitude to the electron Yukawa coupling in the standard model, $h_e \simeq 10^{-6}$ and suggests that the reason why the inflationary potential is so flat is related to the (as yet) unknown reason why the electron Yukawa coupling is very small. Given an intermediate scale breaking $\langle\tilde{P}\rangle \simeq 10^{12}$ GeV, (see the next section) the mass scale of the electron Majorana neutrino would be $M_{N_1} \simeq h_1\langle\tilde{P}\rangle \simeq 10^5$ GeV. The two heavier Majorana neutrino generations are not determined by inflation. If one assumes a hierarchy in the Majorana Yukawa couplings similar to that of the quark and lepton mass spectrum, an interesting light neutrino spectrum can result, with $\Delta m^2_{\mu e} \simeq 10^{-5}\mathrm{eV}^2$ and in certain cases $m_{\nu_e} > m_{\nu_\mu}$. In principle one can also obtain an estimate for the ratio of hot to cold dark matter.

During the initial period of chaotic inflation quantum de-Sitter fluctuations can affect the classical motion of the inflaton. The amplitude of the inflaton decreases exponentially during the de-Sitter phase [9]

$$\tilde{N}^c_1(t) = \tilde{N}^c_1(0)\exp\left[-\frac{h_1}{\sqrt{6\pi}}M_{Pl}t\right]. \tag{5}$$

After a time $\Delta t = H^{-1}$ the amplitude of the inflaton decreases by an amount $\Delta\tilde{N}^c_1 = M^2_{Pl}/(2\pi\tilde{N}^c_1)$, whereas the average amplitude of the quantum fluctuations grows by $\left|\delta\tilde{N}^c_1\right| = H/(2\pi)$. In order for the quantum fluctuations to have negligible influence on the classical evolution $\tilde{N}^c_1(t)$ we need $\tilde{N}^c_1(0) \ll h_1^{-1/3}M_{Pl} \simeq 10^2 M_{Pl}$. In addition the universe must expand greater than 65 e-folds to solve the flatness and horizon problems. This restricts the initial value of the inflaton field to lie in the range $5M_{Pl} \lesssim \tilde{N}^c_1(0) \lesssim 10^2 M_{Pl}$.

3 PQ symmetry breaking

The intermediate scale breaking of PQ symmetry occurs when the singlet scalar field, $\tilde{P}$ receives a vacuum expectation value $\langle\tilde{P}\rangle \simeq 10^{12}$ GeV. Murayama, Suzuki and Yanagida [3] showed that this breaking can be induced by radiative corrections from right-handed

neutrino loops, which drives the mass squared parameter of $\tilde{P}$ negative. The soft supersymmetric breaking terms in the scalar potential involving $\tilde{N}_i^c$,$\tilde{P}$ and $\tilde{P}'$ are given by

$$\begin{aligned} V_{soft} &= m^2_{\tilde{P}} \left|\tilde{P}\right|^2 + m^2_{\tilde{P}'} \left|\tilde{P}'\right|^2 + m^2_{\tilde{N}_i^c} \left|\tilde{N}_i^c\right|^2 + (A_N^{(ij)} h_N^{ij} \tilde{N}_i^c \tilde{L}_j H_u + h.c.) \\ &+ (\tfrac{1}{2} h_i A_i \tilde{N}_i^c \tilde{N}_i^c \tilde{P} + \frac{f}{M_{Pl}} A_f \tilde{P}^3 \tilde{P}' + \frac{g}{M_{Pl}} A_g H_u H_d \tilde{P} \tilde{P}' + h.c.). \end{aligned} \tag{6}$$

The soft scalar masses and trilinear couplings are all apriori unknown mass parameters, but a study of constrained minimal supersymmetry requires them to be $\mathcal{O}(1\ \text{TeV})$[10]. When $m^2_{\tilde{P}} \simeq -m^2_W$, where $m_W \simeq \mathcal{O}(\text{TeV})$ is the electroweak scale, the minimum of the scalar potential

$$V(\tilde{P}) = -m_W^2 \left|\tilde{P}\right|^2 + \frac{f^2}{M_{Pl}^2} \left|\tilde{P}\right|^6 + V_0 \tag{7}$$

occurs for

$$\langle \tilde{P} \rangle = \sqrt{\frac{m_W M_{Pl}}{\sqrt{3} f}} \simeq 10^{12} \text{GeV} \tag{8}$$

where $f \sim 0.01$ and V_0 is the vacuum energy associated with the phase transition. A significantly smaller value of the coupling f would increase the intermediate mass scale and conflict with cosmological axion mass bounds [11]. In addition one finds that to stabilise the scalar potential we need $m^2_{\tilde{P}'} > 0$ and $\langle \tilde{P}' \rangle \simeq \langle \tilde{P} \rangle$ [3]. When the quantum corrections to the soft scalar masses in (6) are included via the one-loop renormalisation group equations, boundary conditions at M_{Planck} determine whether the tree-level result (8) remains valid. In particular for $m^2_{\tilde{P}}$ and $m^2_{\tilde{N}_i^c}$ we have

$$\begin{aligned} \frac{dm^2_{\tilde{P}}}{dt} &= \frac{1}{16\pi^2} \sum_i |h_i|^2 (m^2_{\tilde{P}} + 2m^2_{\tilde{N}_i^c} + |A_i|^2) && (9) \\ \frac{dm^2_{\tilde{N}_i^c}}{dt} &= \frac{1}{16\pi^2} 2 |h_i|^2 (m^2_{\tilde{P}} + 2m^2_{\tilde{N}_i^c} + |A_i|^2). && (10) \end{aligned}$$

Note that in (10) we have not written slepton and Higgs soft mass terms. A complete analysis of all the renormalisation group equations in the MSSM which includes the neutrino masses can have interesting implications. As we noted in the previous section $h_1 \simeq \mathcal{O}(10^{-7})$ and consequently its effect on the renormalisation group running (9) and (10) is negligible when $h_2, h_3 \gg h_1$. This means that the running of $m^2_{\tilde{N}_2^c}$ and $m^2_{\tilde{N}_3^c}$ will be identical to $m^2_{\tilde{P}}$. To ensure that only $m^2_{\tilde{P}}$ goes negative we have to impose the boundary condition $m^2_{\tilde{N}_{2,3}^c} \gtrsim 3m^2_{\tilde{P}}$ at M_{Planck}. Numerical integration of the renormalisation group

equations (9) and (10) with these boundary conditions leads to radiative PQ-symmetry breaking at an intermediate scale, $\langle\tilde{P}\rangle \simeq 10^{12}$GeV.

The radiative corrections indicated by the renormalisation group equations (9) and (10) are evaluated at a temperature $T = 0$ and the quantum fields are assumed to be at their minima. However we need to include corrections arising from the inflationary period and consider possible thermal effects. During the inflationary epoch the inflaton field sits far from its minimum with a value $\tilde{N}_1^c(0) \gtrsim M_{Pl}$. As noted earlier the inflaton can induce masses to any other scalar fields that it couples to in the scalar potential (3). While the Higgs and slepton fields receive an effective mass $\gtrsim H$, the coupling $h_1^2 \left|\tilde{N}_1^c\right|^2 \left|\tilde{P}\right|^2$ induces an effective mass $h_1\langle\tilde{N}_1^c(t)\rangle$ for $\tilde{P}$. This mass will dominate any radiative corrections until the inflaton field $\tilde{N}_1^c(t)$ settles to its minimum after inflation ends.

The finite temperature corrections associated with the potential (7) have been previously discussed in the context of intermediate scale breaking in superstring models [12]. The finite temperature potential for $m_W \ll T \ll M_I$ and excluding the region $T \sim \left|\tilde{P}\right|$ is given by

$$
\begin{aligned}
V(\left|\tilde{P}\right|, T) &\simeq -m_W^2 \left|\tilde{P}\right|^2 + \frac{\pi^2}{90} T^4 \qquad (T \ll \left|\tilde{P}\right| < M_I) &\quad (11)\\
&\simeq \frac{h_3^2}{24} T^2 \left|\tilde{P}\right|^2 \qquad (T \gg \left|\tilde{P}\right|) &\quad (12)
\end{aligned}
$$

where M_I is the intermediate breaking scale. Since the third generation Majorana neutrino Yukawa coupling $h_3 \simeq 1$, the finite temperature potential has a local minimum at $\langle\tilde{P}\rangle = 0$ which disappears when $T \lesssim m_W \simeq 10^3$GeV. The problem is that when the universe is supercooled at the end of inflation, the inflaton induced $\tilde{P}$ mass ($h_1\langle\tilde{N}_1^c(t)\rangle$) still dominates the radiative corrections ($\langle\tilde{N}_1^c\rangle \simeq \frac{1}{3}M_{Pl}$) and so $\langle\tilde{P}\rangle \simeq 0$. When the universe reheats to a temperature $T_{RH} \simeq 10^4$GeV the scalar field $\tilde{P}$ is still trapped at the origin with a barrier of height $\sim T^4$. Eventually when $T \simeq m_W$ the barrier disappears and then $\langle\tilde{P}\rangle \simeq M_I \simeq 10^{12}$GeV. Electroweak baryogenesis will then be the only possibility for generating a baryon asymmetry.

However, in general there are many flat directions in supersymmetric theories and it is very likely that a flat-direction field, η has an amplitude $\mathcal{O}(M_{Pl})$. This can occur via quantum de-Sitter fluctuations along the F and D-flat directions or it can be driven to an $\mathcal{O}(M_{Pl})$ local minimum by non-renormalisable Kahler potential couplings during the initial inflationary epoch (see Dine et al [5]). If we assume this is the case then as the right-handed electron sneutrino continues to roll towards its minimum, there will

be a point where the flat-direction field η dominates the potential energy density with $\eta(0) \simeq M_{Pl}$ and $V(\eta) = \frac{1}{2}m_W^2\eta^2$. A second period of chaotic inflation will then commence, which accelerates the damping of the $\tilde{N}_1^c$ oscillations and supercools the universe again. Eventually the $\tilde{P}$ soft term will dominate the inflaton induced $\tilde{P}$ mass (temperature effects are negligible) and generate an intermediate scale ($\langle\tilde{P}\rangle \simeq M_I$). We can neglect the quantum de-Sitter fluctuations during this second period of inflation because $\sqrt{\langle\chi^2\rangle} \lesssim m_W$. The pseudo-Nambu-Goldstone boson resulting from the spontaneous symmetry breakdown will be the invisible axion. The right-handed electron sneutrino and electron Majorana neutrino will then become massive with $M_{N_1} \simeq 10^5$ GeV. In addition the MSSM Higgs mass term ($\mu\hat{H}_u\hat{H}_d$) is generated with $\mu \simeq \mathcal{O}(m_W)$.

The number of e-foldings N, produced during the intermediate scale inflation depends on the initial value of the flat-direction field and is given by $N \simeq 2\pi\eta(0)^2/M_{Pl}^2$. Typically we expect $\eta(0) \lesssim 2M_{Pl}$ to avoid fine-tuning problems and so $N \lesssim 25$. This amount of inflation is not enough to expand different θ_i axion domains beyond our present day horizon, so cosmic axion strings will be present (although the strings will be diluted). However, even if axion strings occur and lead to domain walls at the QCD transition temperature it is not clear that this leads to any cosmological problems [13].

The adiabatic density perturbations produced by this second period of inflation will be negligible because $\delta\rho/\rho \simeq m_W/M_{Pl} \simeq 10^{-16}$. In order that they become irrelevant for galaxy formation and not destroy the density perturbations produced by $\tilde{N}_1^c$, one requires that the density perturbations re-enter the horizon for time scales irrelevant to the growth of large scale cosmological density perturbations. This requires that the number of e-foldings $N \leq 30$ for $T_{RH} \simeq 10^6$GeV [7] which is satisfied for $\eta(0) \simeq 2M_{Pl}$. Note also that quantum fluctuations of the axion field can produce isothermal density perturbations, but these will be negligible because the Hubble constant during this second inflationary epoch is $H \simeq m_W$ [11].

When the second period of inflation ends the universe will be reheated to a temperature $T_{RH} \simeq g_\star^{-1/4}\sqrt{\Gamma_\eta M_{Pl}} \simeq 10^6$GeV, where $g_\star \simeq 280$ and $\Gamma_\eta \simeq h_Y^2 m_W/(4\pi)$ for an inflaton with a Yukawa-type coupling $h_Y \simeq 10^{-4}$. Since $\tilde{P}$ is sitting (up to quantum fluctuations $\lesssim \mathcal{O}(m_W)$) at the global minimum $\langle\tilde{P}\rangle \simeq M_I$ with a potential depth $m_W^2 M_I^2 \simeq (10^7\text{GeV})^4$, finite temperature corrections do not destroy this minimum, even though a local minimum exists at the origin. Note that the reheat temperature is sufficiently low to avoid the gravitino problem [14]. In addition the reheat temperature is sufficiently high that right-handed electron neutrinos (N_1) are regenerated because $M_{N_1} \simeq 10^5$ GeV. When $T \simeq M_{N_1}$ a lepton asymmetry will be generated by out of equi-

librium CP-violating decays of N_1, provided that the neutrino Dirac Yukawa couplings are complex and $\left|h_N^{1j}\right| \simeq 10^{-6}$. This lepton asymmetry will be reprocessed into a baryon asymmetry by the usual electroweak anomaly.

4 Conclusion

We have described how one can obtain chaotic inflation with a radiatively generated intermediate mass scale in a simple phenomenological extension of the MSSM. We build on but significantly modify the approach of Murayama et al [1]. An initial period of inflation, driven by a quartic potential associated with the right-handed electron sneutrino solves the usual horizon and flatness problems of the universe. Density perturbations, $\delta\rho/\rho \simeq 10^{-5}$ are generated when the electron Majorana neutrino Yukawa coupling is $\mathcal{O}(10^{-7})$. While technically natural, this coupling is similiar in magnitude to the electron Dirac Yukawa coupling in the MSSM. Radiative corrections from right-handed neutrino loops will break $\mathrm{U}(1)_{PQ}$ at an intermediate scale (10^{12}GeV) when the universe cools to a temperature $T \lesssim 10^3$GeV. This implies an electron Majorana neutrino mass $M_{N_1} \simeq 10^5$GeV and suggests that there is a hierarchy in the Majorana neutrino spectrum related to the quark and lepton spectrum. A baryon asymmetry can only be generated by invoking the non-perturbative processes of electroweak baryogenesis.

However it is possible that a second stage of inflation can occur at an intermediate scale with some flat-direction field, η. This second inflationary epoch accelerates the damping of the electron sneutrino amplitude and exponentially cools the universe again, causing the radiative corrections to dominate and spontaneously break $\mathrm{U}(1)_{PQ}$. When inflation ends the universe can be reheated to a temperature $T_{RH} \simeq 10^6$GeV which is sufficiently low to prevent restoring PQ symmetry. In addition the right-handed electron neutrinos are regenerated and eventually decay when $T \simeq M_{N_1}$. The resulting lepton asymmetry is reprocessed by the electroweak anomaly into a baryon asymmetry.

The model we have constructed is an attempt to amalgamate current cosmological ideas with the well tested phenomenology of the MSSM. Ultimately one would like motivation from a more fundamental theory, but we hope that the effective Lagrangian we have considered can shed some light in this direction.

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